Couplings of quarks in the Partially Aligned 2HDM with a four-zero texture Yukawa matrix

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Abstract

The Two Higgs Doublets Model (2HDM) has provided a very useful way to describe a minimal extension of the scalar sector of the Standard Model. In this work, it is shown a scheme that we call Partial Aligned Two Higgs Doublet Model (PA-2HDM) which allows a description of the distinct versions of the 2HDM in a simple way, including those with flavor symmetries. In addition, it is shown a method to diagonalize Yukawa matrices of four-zero texture coming from the 2HDM-III. We provide some phenomenological applications in order to show the model's predictive power.

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The main problem in the flavor physics beyond the Standard Model (SM) [1] is to control the presence of the Flavor Changing Neutral Currents (FCNC) that in experiments have been observed to be highly suppressed. Almost all models that describe physics in energy regions greater than the electroweak scale, have contributions with FCNC at tree level, unless some symmetry is introduced on the scalar sector to suppress them. One of the most important extensions of the SM is the Two Higgs Doublet Model (2HDM), due to its wide variety of dynamical features and the fact that it can represent a low-energy limit of general models like the Minimal Supersymmetric Standard Model. There are some generalizations of the 2HDMs of type I, II, X and Y (2HDM-I, 2HDM-II, 2HDM-X and 2HDM-Y) [2], as well as the 2HDM-III with flavor symmetries that require a four texture in the Yukawa matrix [3] and Lepton Flavor Violating (LFV) introduced as a deviation from Model II Yukawa interaction [4, 5]. The type-X (type-Y) 2HDM is referred to as the type-IV (type-III) 2HDM in Ref.[6] and the type-I' (type-II') 2HDM in Ref. [7, 8]. Sometimes, the most general 2HDM, in which each fermion couples to both Higgs doublet fields, is called the type III 2HDM [9]. From a phenomenological point of view, the Cheng-Sher *ansatz* [10] has been very useful to describe the phenomenological content of the Yukawa matrix and the salient feature of the hierarchy of quark masses. Through the Yukawa textures [11, 12] it is possible to build a matrix that preserves the expected Yukawa couplings that depend on the fermion masses. One unavoidable problem is the great number of free parameters that emerge as a consequence of introducing a new Higgs doublet. In order to reduce the number of free parameters of the model some restrictions have been imposed on the entries of the Yukawa matrices through discrete symmetries or phenomenological assumptions. The Yukawa Lagrangian for the quark fields is given by

$$\mathcal{L}_Y = Y_1^u \bar{Q}_L \tilde{\Phi}_1 u_R + Y_2^u \bar{Q}_L \tilde{\Phi}_2 u_R + Y_1^d \bar{Q}_L \Phi_1 d_R + Y_2^d \bar{Q}_L \Phi_2 d_R \tag{1}$$

where $\Phi_{1,2} = (\phi_{1,2}^+, \phi_{1,2}^0)^T$ denotes the Higgs doublets, $\tilde{\Phi}_{1,2} = i\sigma_2 \Phi_{1,2}^*$ and $Y^{u,d}$ are the Yukawa matrices.

The above Lagrangian (1) has a great deal of free parameters associated with the Yukawa interaction and five scalar bosons, two of them charged ($H^\pm$) and one of the neutral ones is a pseudoscalar (A^0). The mechanism through which the FCNC are controlled defines the version of the model and a different phenomenology that can be contrasted with the experiment. One successful version where the Yukawa couplings depend on the hierarchy of masses is the one where the mass matrix has a four-zero texture form. This matrix is based on the phenomenological observation that the off-diagonal elements must be small in order to dim the interactions that violate flavor as the experimental results show. Although the phenomenology of Yukawa couplings

constrains the hierarchy of the mass matrix entries, it is not enough to determine the strength of the interaction with scalars. Another assumption on the Yukawa matrix is related to the additional Higgs doublet. In versions I and II it is introduced a discrete symmetry on the Higgs doublets, fulfilled by the scalar potential, that leads to the vanishing of most of the free parameters. However, version III, having a richer phenomenology, requires a slightly more general scheme.

There is a close relation between the flavor space and the mass matrix, which in general can be written as

$$M_f = \frac{1}{\sqrt{2}}(v_1 Y_1^f + v_2 Y_2^f). \tag{2}$$

Inspired by the fact that in the Higgs basis the information of the mass matrix is contained in the first Yukawa matrix and the SM couplings are proportional to the fermion masses, the interactions with scalars in a general 2HDM can be modeled by imposing a specific form on the second Yukawa matrix as a mass matrix transformed in the flavor space. In this paper it is utilized a particular case of this model [13] that can describe different versions of the 2HDM by using properties of the flavor space through a simple principle. We introduce the concept Partially Aligned (PA) Yukawa Matrix according to two criteria: a) a new transformation for the first Yukawa matrix in the flavor space $SU_F(3)$ and b) the control of FCNC induced by this transformation, using as a criterion the Cheng-Sher *ansatz* [10]. By following these ideas, the concept of Partially Aligned (PA) will be defined by a new transformation which enables us to write the matrix of couplings as a bi-unitary transformed mass matrix, namely

$$Y_2^f = \frac{1}{v} A_L^f M_f A_R^f, \tag{3}$$

where A_L^f and A_R^f with $f = u, d, \ell$, are diagonal $SU_F(3)$ matrices that concentrate the dynamical information about extended scalar interactions and M_f contains the properties of the hierarchy of the quark masses and the mixing of the CKM matrix, whose form is determined by a more fundamental theory. As usual, we have combined the VEVs of the doublet Higgs fields through the relation $v^2 = v_1^2 + v_2^2$. In the PA-2HDM the aligned model [14, 15] can be cast with $A_L^f = A_R^f \sim \lambda_0$, where λ_0 is the matrix proportional to a unit matrix in $SU_F(3)$. Details about these formulations are given elsewhere. As mentioned above, the several versions of the 2HDM can be generated by choosing suitable matrices (see table I). There is no physical restriction on the structure of the mass matrix beyond the fact that the quark masses of different families differ by several orders of magnitude.

	A_L^u	A_R^u	A_L^d	A_R^d
I	$\sqrt{\frac{3m_W}{v}}\lambda_0$	$\sqrt{\frac{3m_W}{v}}\lambda_0$	$\sqrt{\frac{3m_W}{v}}\lambda_0$	$\sqrt{\frac{3m_W}{v}}\lambda_0$
II	$\sqrt{\frac{3m_W}{v}}\lambda_0$	$\sqrt{\frac{3m_W}{v}}\lambda_0$	$0_{3\times3}$	$0_{3\times3}$
III-IV	$\sum_{a=0,3,8} C_a^u\lambda_a$	$\left(\sum_{a=0,3,8} \tilde{C}_a^u\lambda_a\right)^\dagger$	$\sum_{a=0,3,8} C_a^d\lambda_a$	$\left(\sum_{a=0,3,8} \tilde{C}_a^d\lambda_a\right)^\dagger$
A2HDM	$C_0^u\lambda_0$	$\tilde{C}_0^{u*}\lambda_0$	$C_0^d\lambda_0$	$\tilde{C}_0^{d*}\lambda_0$

TABLE I. Matrices that reproduce several versions of the Yukawa interactions for the 2HDM in terms of $SU_F(3)$ generators. The $C's$ parameters are complex coefficients and they are proportional to the parameters $\widetilde{\chi}_{ij}^f$ defined in Eq.(6).

On the other hand, the PA-2HDM will induce Higgs boson FCNC through the following term

$$\widetilde{Y}_2^f = \frac{1}{v}\widetilde{A}_L^f \bar{M}_f \widetilde{A}_R^f, \tag{4}$$

where $\widetilde{A}_{L,R}^f = U_{L,R}^{f\dagger} A_{L,R}^f U_{L,R}^f$, $\bar{M}_f = \mathrm{Diag}[m_{f1}, m_{f2}, m_{f3}]$ and $U_{L,R}^f$ are the matrices that diagonalize the mass matrix M_f. So, the contribution to fermion-fermion-Higgs bosons couplings is given by:

$$(\widetilde{Y}_2^f)_{ij} = \frac{1}{v}\left(m_{f1}(\widetilde{A}_L^f)_{i1}(\widetilde{A}_R^f)_{1j} + m_{f2}(\widetilde{A}_L^f)_{i2}(\widetilde{A}_R^f)_{2j} + m_{f3}(\widetilde{A}_L^f)_{i3}(\widetilde{A}_R^f)_{3j}\right). \tag{5}$$

In order to control the FCNC induced by the model, we employ the Cheng-Sher *ansatz* [10] in the following way:

$$(\widetilde{Y}_2^{CS,f})_{ij} = \frac{\sqrt{m_{fi}m_{fj}}}{v}\widetilde{\chi}_{ij}^f, \tag{6}$$

then, from Eq.(4) and Eq.(6) the FCNC will be controlled by:

$$\left|m_{f1}(\widetilde{A}_L^f)_{i1}(\widetilde{A}_R^f)_{1j} + m_{f2}(\widetilde{A}_L^f)_{i2}(\widetilde{A}_R^f)_{2j} + m_{f3}(\widetilde{A}_L^f)_{i3}(\widetilde{A}_R^f)_{3j}\right| \leq \sqrt{m_{fi}m_{fj}}\left|\widetilde{\chi}_{ij}^f\right|. \tag{7}$$

The advantage of this criterion is that we can use previous studies of the experimental constraints imposed on the free parameters of Yukawa texture [3, 16, 18–20]. Moreover, by definition, the eigenvalues of the mass matrix are the masses of fermions, *i.e.*, they must be real and non-negative. A hermitian matrix guarantees that the masses are real, however, the non-negativity condition for the eigenvalues is not fulfilled by any hermitian matrix. Actually, the four-zero texture matrix has at least one negative eigenvalue. This drawback is solved by considering that upon diagonalization the masses are the square root of the eigenvalues of $H_f = M_f M_f^\dagger$. This assumption is appropriate to determine the phenomenological couplings in the Yukawa sector as the Cheng-Sher *ansatz*, albeit, it leaves out other possible parameterizations emerging from a rather general method.

In the following, instead of making assumptions on the nature of the eigenvalues of M_f, we look for the properties of free parameters of the mass matrix in the flavor basis in order to generate real and non-negative eigenvalues of M_f. In general, the bi-unitary transformation is given by $M_f = U_L^f M_f U_R^{f\dagger}$, where U_L^f and U_R^f represent the unitary transformations that diagonalize $M_f M_f^\dagger$ and $M_f^\dagger M_f$, respectively. If M_f were hermitian then U_L^f and U_R^f would be equal. In what follows we shall restrict ourselves to this case. The most suitable mass matrix structure that describes the properties of the couplings of the Yukawa sector is the four-zero texture Yukawa matrix, which can be written as [3, 16]

$$M_f = \begin{pmatrix} 0 & C_f & 0 \\ C_f^* & \widetilde{B}_f & B_f \\ 0 & B_f^* & A_f \end{pmatrix}. \tag{8}$$

The form and hierarchy of the free parameters are conserved, although, for the sake of generalization, we slightly break the hermitian condition to allow for new possible effects coming from relative phases between diagonal elements A_f and $\tilde{B}_f$. It is worth mentioning that these phases must obey the physical condition imposed on the eigenvalues of the mass matrix,

$$A_f = |A_f| \cdot e^{i\theta_{Af}}, \tag{9}$$

$$\widetilde{B}_f = |\widetilde{B}_f| \cdot e^{i\theta_{\widetilde{B}f}}. \tag{10}$$

The hermitian matrix H_f is given by

$$H_f = \begin{pmatrix} |C_f|^2 & \widetilde{B}_f^* C_f & B_f C_f \\ \widetilde{B}_f C_f^* & |C_f|^2 + |\widetilde{B}_f|^2 + |B_f|^2 & \widetilde{B}_f B_f + A_f^* B_f \\ B_f^* C_f^* & \widetilde{B}_f^* B_f^* + A_f B_f^* & |A_f|^2 + |B_f|^2 \end{pmatrix}. \tag{11}$$

As usual, we can extract the phases of the non-diagonal elements with a transformation $H_f = P_f^\dagger \overline{H}_f P_f$ where $P_f = \mathrm{Diag}[1, e^{-i(\theta_{\widetilde{B}f} - \theta_{Cf})}, e^{i(\theta_{Bf} + \theta_{Cf})}]$. It is important to highlight the unitary relation between H_r and the diagonal matrix $M_f^2 = \mathrm{Diag}[m_{f1}^2, m_{f2}^2, m_{f3}^2]$ that leads to the following system of equations:

$$|A_f|^2 + 2|B_f|^2 + 2|C_f|^2 + |\widetilde{B}_f|^2 = m_{f1}^2 + m_{f2}^2 + m_{f3}^2, \tag{12}$$

$$\begin{aligned}&(|B_f|^2 + |C_f|^2)^2 + |A_f|^2|\widetilde{B}_f|^2 + 2|A_f|^2|C_f|^2 - 2\cos(\theta_{Af} + \theta_{\widetilde{B}f})|A_f||\widetilde{B}_f||B_f|^2 \\ &= m_{f1}^2 m_{f2}^2 + m_{f1}^2 m_{f3}^2 + m_{f2}^2 m_{f3}^2,\end{aligned} \tag{13}$$

$$|A_f|^2|C_f|^4 = m_{f1}^2 m_{f2}^2 m_{f3}^2. \tag{14}$$

Thus, the problem is reduced to solving the system of equations, given by Eqs.(12-14), which has 16 possible solutions, though most of them unphysical. However, by adopting a simple criterion we can to simplify this system of equations. As mentioned above, we are only interested in those solutions which reproduce real and non-negative eigenvalues. To achieve this goal we establish the following three conditions that ensure physical properties of the solutions:

- The free complex parameters A_f, B_f, C_f and $\widetilde{B}_f$ are a function of the masses and satisfy the invariant equations (12-14).

- The eigenvalues of M_f, namely, the fermion masses, must be real and non-negative.

- The eigenvalues of H_f must obey the hierarchy of the quark masses as experimentally observed, *i.e.*, $m_{f3} > m_{f2} > m_{f1}$ [17].

A more simplified system of equations is thus obtained by factorization due to the chosen phases which must fulfill the above conditions:

$$|A_f| + (-1)^m|\widetilde{B}_f| = m_{f1} - m_{f2} + m_{f3}, \tag{15}$$

$$|B_f|^2 - (-1)^m|A_f||\widetilde{B}_f| + |C_f|^2 = m_{f1}\, m_{f2} - m_{f1}\, m_{f3} + m_{f2}\, m_{f3}, \tag{16}$$

$$|A_f||C_f|^2 = m_{f1}\, m_{f2}\, m_{f3}. \tag{17}$$

with m integer. The solution is given by

$$|B_f| = \sqrt{\left(1 - \frac{m_{f1}}{|A_f|}\right)(|A_f| + m_{f2})(m_{f3} - |A_f|)}, \tag{18}$$

$$|\widetilde{B}_f| = (-1)^m(m_{f1} - m_{f2} + m_{f3} - |A_f|), \tag{19}$$

$$|C_f| = \sqrt{\frac{m_{f1}\, m_{f2}\, m_{f3}}{|A_f|}}. \tag{20}$$

Therefore, the value of parameter $|A_f|$ depends on the parity of m: for m even we have $m_{f1} \leq |A_f| \leq m_{f1} - m_{f2} + m_{f3}$, whereas m odd leads to $m_{f1} - m_{f2} + m_{f3} \leq |A_f| \leq m_{f3}$. We shall assume a linear behavior of $|A_f|$ in terms of the parameters $0 \leq \beta_i^f \leq 1$ $(i = 1, 2)$, so that one can write $|A_f| = m_{f1}\left(1 + \beta_1^f \frac{m_{f3}-m_{f2}}{m_{f1}}\right)$, for m even, and $|A_f| = m_{f3}\left(1 - \beta_2^f \frac{m_{f2}-m_{f1}}{m_{f3}}\right)$, for m odd. The idea is then to expand $|A_f|$ in terms of $z = \frac{m_{f3}-m_{f2}}{m_{f1}}$, for m even, and $z = \frac{m_{f2}-m_{f1}}{m_{f3}}$, for m odd. Considering now a four-zero texture form for the mass matrix M_f and choosing

$A_L^f = \mathrm{Diag}\left[1, \frac{d_2^f}{c_2^f}, \frac{b_2^{f*}}{c_2^f}\right]$ and $A_R^f = \mathrm{Diag}\left[\frac{|c_2^f|^2}{d_2^f}, c_2^f, \frac{b_2^f c_2^f}{d_2^f}\right]$, we thus have

$$Y_2^f = \begin{pmatrix} 0 & c_2^f\, C_f & 0 \\ c_2^{f*}\, C_f^* & d_2^f \widetilde{B}_f & b_2^f B_f \\ 0 & b_2^f B_2^{f*} & a_2^f A_f \end{pmatrix}, \tag{21}$$

where $a_2^f = \frac{|b_2^f|^2}{d_2^f}$. The Yukawa matrix preserves the four-zero texture form. Thus, for m odd, $|A_f| = m_{f3}\left(1 - \beta_2^f \frac{m_{f2}-m_{f1}}{m_{f3}}\right)$, one can reproduce the parametrization of a four-zero texture Yukawa matrix given in Ref.[3, 16]. So, the Cheng-Sher *ansatz* from Eq.(6) can be reproduced in the limit $m_{f1} << m_{f2} << m_{f3}$, and the parameters $\widetilde{\chi}_{ij}^f$ can be written in terms of the entries of A_L^f and A_R^f matrices,

$$\widetilde{\chi}_{11}^f = \left[d_2^f - (c_2^{f*} e^{i\phi_{cf}} + c_2^f e^{-i\phi_{cf}}),\right] \eta^f + \left[a_2^f + d_2^f - (b_2^{f*} e^{i\theta_{bf}} + b_2^f e^{-i\theta_{bf}})\right] \beta_2^f, \tag{22}$$

$$\widetilde{\chi}_{12}^f = c_2^f e^{-i\theta_{cf}} - d_2^f - \eta^f \left[a_2^f + d_2^f - (b_2^{f*} e^{i\theta_{Bf}} + b_2^f e^{-i\theta_{Bf}})\right] \beta_2^f, \tag{23}$$

$$\widetilde{\chi}_{13}^f = (a_2^f - b_2^f e^{-i\theta_{Bf}})\eta^f \sqrt{\beta_2^f}, \tag{24}$$

$$\widetilde{\chi}_{22}^f = d_2^f \eta^f + \left[a_2^f + d_2^f - (b_2^{f*} e^{i\theta_{Bf}} + b_2^f e^{-i\theta_{Bf}})\right] \beta_2^f, \tag{25}$$

$$\widetilde{\chi}_{23}^f = (b_2^f e^{-i\theta_{Bf}} - a_2^f)\sqrt{\beta_2^f}, \tag{26}$$

$$\widetilde{\chi}_{33}^f = a_2^f, \tag{27}$$

where $\eta^f = \lambda_2^f / m_2^f$, with $m_2^f = |\lambda_2^f|$, and $0 \leq \beta_2^f \leq 1$. The above equations can be inverted to estimate the entries of matrices A_L and A_R. This case has been studied previously and the same constraints in the parameters χ_{ij} can be imposed. In particular, we find that $\chi_{ij}^q = O(1)$ are allowed [3, 16, 18–20]. In figure 1 are shown the values of the Yukawa matrix entries when $|\widetilde{\chi}_{ij}| \sim 1$ and the phases are taken to be zero. The sudden fall and rise of $|(Y_2^f)_{12}|$ in the upper right plot of Fig.1 stems from a sign change in the Yukawa coupling value. This case represents a special approximation when $\mathrm{Arg}(C_a^f) = \mathrm{Arg}(\widetilde{C}_a^f)$ (see table I).

On the other hand, for m even $|A_f| = m_{f1}\left(1 + \beta_1^f \frac{m_{f3}-m_{f2}}{m_{f1}}\right)$, in which case, one can see from Eq. (21) that the Yukawa coupling form changes and therefore its parametrization. Although the analytical expressions of Yukawa texture for m-even are larger than for the m-odd case, in general, we get that $(Y_2^f)_{m-even} \propto (Y_2^f)_{m-odd}$. We shall discuss in in detail the phenomenology of this scenario in a forthcoming paper, however, for practical reasons, this case is here analyzed numerically. For the quark sector it is possible to estimate the value of β_1^u, β_1^d by using the experimental information of the CKM matrix. Our analysis gives $\beta_1^u = \beta_1^d \sim 0.9985$, to be compared

with the Yukawa couplings including the Cheng-Sher parametrization. For the up-type quark sector we have

$$\frac{|(Y_2^{even,u})_{11}|}{|(Y_2^{CS,u})_{11}|} \sim 13.08; \qquad \frac{|(Y_2^{even,u})_{12}|}{|(Y_2^{CS,u})_{12}|} \sim 6.85; \qquad \frac{|(Y_2^{even,u})_{13}|}{|(Y_2^{CS,u})_{13}|} \sim 8.78, \tag{28}$$

$$\frac{|(Y_2^{even,u})_{22}|}{|(Y_2^{CS,u})_{22}|} \sim 3.56; \qquad \frac{|(Y_2^{even,u})_{23}|}{|(Y_2^{CS,u})_{23}|} \sim 5.00; \qquad \frac{|(Y_2^{even,u})_{33}|}{|(Y_2^{CS,u})_{33}|} \sim 0.80. \tag{29}$$

In the down-type quark sector,

$$\frac{|(Y_2^{even,d})_{11}|}{|(Y_2^{CS,d})_{11}|} \sim 8.38; \qquad \frac{|(Y_2^{even,d})_{12}|}{|(Y_2^{CS,d})_{12}|} \sim 3.69; \qquad \frac{|(Y_2^{even,d})_{13}|}{|(Y_2^{CS,d})_{13}|} \sim 5.13, \tag{30}$$

$$\frac{|(Y_2^{even,d})_{22}|}{|(Y_2^{CS,d})_{22}|} \sim 1.57; \qquad \frac{|(Y_2^{even,d})_{23}|}{|(Y_2^{CS,d})_{23}|} \sim 2.64; \qquad \frac{|(Y_2^{even,d})_{33}|}{|(Y_2^{CS,d})_{33}|} \sim 0.69. \tag{31}$$

and FCNC are under control. For leptons (ℓ) we obtain for all cases that

$$\frac{|(Y_2^{even,\ell})_{ij}|}{|(Y_2^{CS,\ell})_{ij}|} \sim O(1). \tag{32}$$

Based on these results and following Eqs. (5-7), one can obtain for all fermions

$$\frac{|\widetilde{\chi}_{ij}^{even,f}|}{|\widetilde{\chi}_{ij}^{CS,f}|} = \frac{|(Y_2^{even,f})_{ij}|}{|(Y_2^{CS,f})_{ij}|}. \tag{33}$$

We have thus managed to implement all experimental constraints found previously [3, 16, 18–20]. In this report we find strong constraints for the free parameters $\widetilde{\chi}_{ij}^{even,f}$. Following References [3, 16], we can obtain the constraint $|\widetilde{\chi}_{12}^{even,\ell}| \leq 5\times10^{-1}$ from $\mu^- - e^-$ conversion, $|\widetilde{\chi}_{13}^{even,\ell}| = |\widetilde{\chi}_{23}^{even,\ell}| \leq 10^{-2}$ from radiative decay $\mu^+ \to e^+\gamma$, and $|\widetilde{\chi}_{23}^{even,d}| \leq 0.2$ from the contribution to the decay $b \to s\gamma$ measurements. In addition, without loss of generality, we can implement all previous studies given in references [3, 16, 18–20] and we can validate the PA-2HDM as a framework phenomenologically viable, as well as the corresponding predictions.

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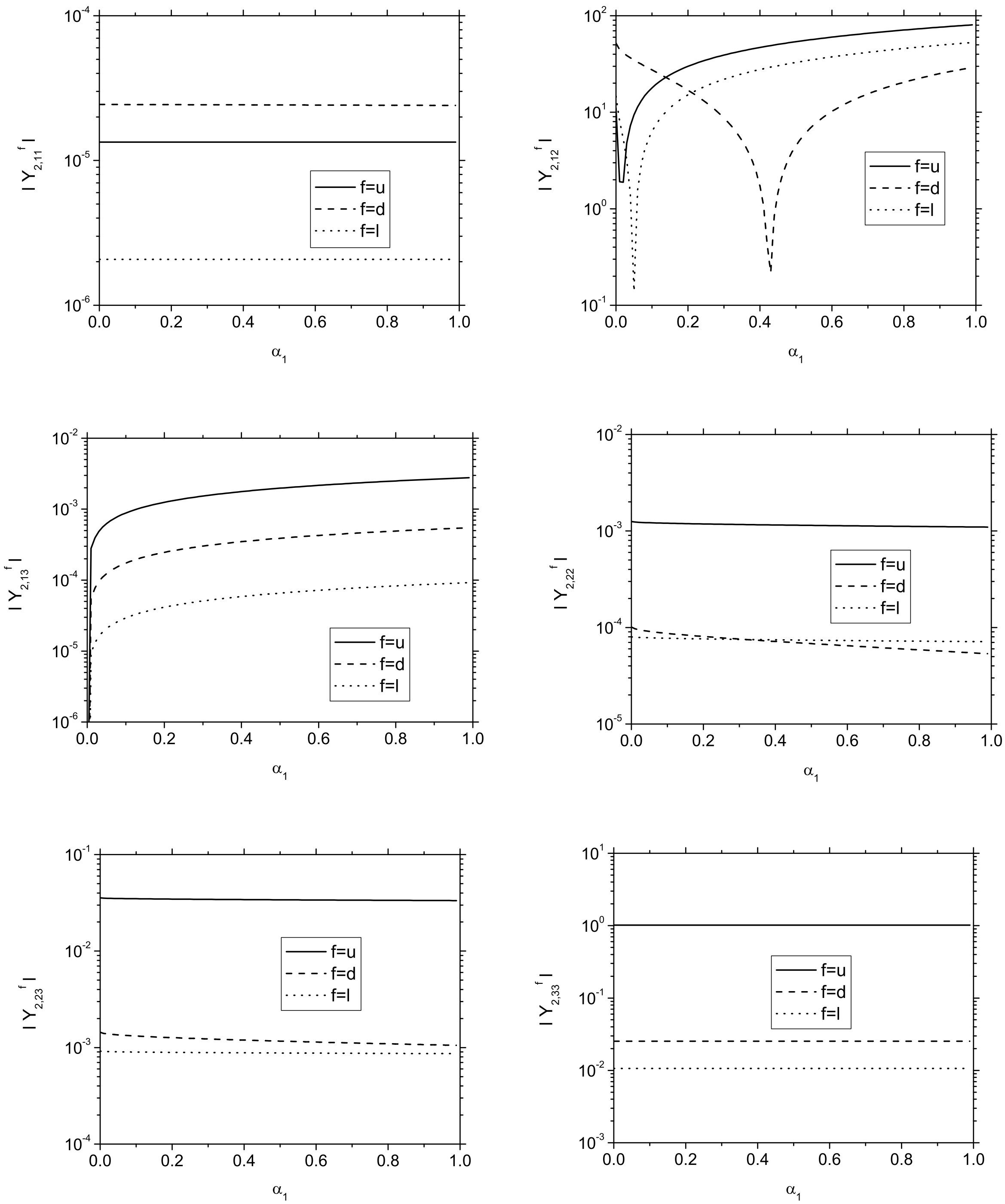

FIG. 1. Magnitude of $|Y^f_{2,ij}|$ in the limit when the Cheng-Sher couplings are $|\widetilde{\chi}_{ij}| \sim 1$ and the phases are taken to be zero